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Heavy quark spin selection rule and the properties of the $X(3872)$

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Abstract

The properties of the resonance $X(3872)$ are discussed under the assumption that this resonance is dominantly a 'molecular' $J^{PC} = 1^{++}$ state of neutral D and D^* mesons. It is argued that in these properties should dominate the states with the total spin of the charmed quark-antiquark pair equal to one. As a practical application of this observation the ratio of the rates of the decays $X \to \pi^0 \chi_{cJ}$ for different J is predicted. It is also pointed out that the total rate of these decays is likely to be comparable to that of the observed transitions $X \to \pi^+ \pi^- J\psi$ and $X \to \pi^+ \pi^- \pi^0 J\psi$. The decays of the X into light hadrons and its production in hadronic processes are also briefly discussed.

The most recent observation[1] of the decay $X(3872) \to \pi^+ \pi^- \pi^0 J/\psi$ with the rate comparable to that of the previously observed[2, 3, 4, 5] process $X(3872) \to \pi^+ \pi^- J/\psi$ unambiguously illustrates (through the obvious G parity reasoning) that the resonance $X(3872)$ has no definite isospin, and thus rules out an interpretation of this resonance as a state of charmonium. Furthermore, the spectra of the di- and tri- pion invariant masses strongly suggest that the underlying processes in the observed decays are $X \to \omega J/\psi$ and $X \to \rho^0 J/\psi$. The extreme proximity to the kinematical boundary in the latter transitions favors interpreting them as S wave processes, which uniquely points to the $J^{PC} = 1^{++}$ assignment for the $X(3872)$, given the conspicuous absence of its decays into pairs of pseudoscalar D mesons. Also the fact that the mass of $X(3872)$ is within approximately $1\,MeV$ or less from the sum of the masses of D^0 and D^{*0} makes it more than natural to interpret the discussed resonance as dominantly a C-parity even S-wave near- threshold state of the neutral meson-antimeson pair: $(D^0 \bar{D}^{*0} + \bar{D}^0 D^{*0})/\sqrt{2}$, much in the same way as the deuteron is dominantly a state made of a proton and neutron. Such 'molecular' systems made of charmed mesons were suggested[6] and discussed[7, 8, 9] long ago, and this interpretation was considered as a favored option for the $X(3872)$[10, 11, 12, 13] ever since its first observation[2]. In particular the observed rate of the decay $X(3872) \to \pi^+ \pi^- \pi^0 J/\psi$ is in a very reasonable agreement with the prediction[14] from a specific model of the internal 'molecular' dynamics of the $X(3872)$.

Clearly, the notion of a 'molecular' state can refer only to the peripheral part of the internal wave function of X, i.e at long distances beyond the range of the strong interaction. In particular, if the mass of the $X(3872)$ is below the $D^0 \bar{D}^{*0}$ threshold by the gap w, the peripheral part of its wave function is that of a free motion of heavy mesons at the characteristic distances set by the scale $1/\sqrt{m_D\, w}$, which is larger than about $5\,fm$ for $w < 1\,MeV$. At shorter distances, i.e. within the range of the strong interaction, the mesons overlap and the "core" of the wave function is determined by multi-body dynamics of heavy and light quarks and gluons. Unlike the peripheral part, which is described by just one state of the meson pair, the properties of the core are determined by a significantly larger number of states in the Fock decomposition

$$\psi_X = a_0\, \psi_0 + \sum_i a_i\, \psi_i \ , \tag{1}$$

where ψ_0 is the state $(D^0 \bar{D}^{*0} + \bar{D}^0 D^{*0})/\sqrt{2}$, while ψ_i refer to 'other' hadronic states. The notion of X being a 'molecular' system is helpful inasmuch as the probability weight $|a_0|^2$ of

the meson component ψ_0 makes a large portion of the total normalization. In particular the model of Ref.[14] includes S and D wave states of the neutral and charged charmed meson pairs as well as the channels $\rho J/\psi$ and $\omega J/\psi$, and estimates the weight factor of the S wave $(D^0\bar{D}^{*0} + \bar{D}^0 D^{*0})$ component as 70-80% at $w \approx 1\,MeV$.

It can be noted however, that although the peripheral 'molecular' component ψ_0 has the largest probability, only a limited scope of properties of the X can be understood using only this component, such as (yet unobserved) decays $X \to D^0\,\bar{D}^0\,\pi^0$ and $X \to D^0\,\bar{D}^0\,\gamma$[12], for which the underlying processes is the decay of individual D^* $(\bar{D}^*)$ meson. The majority of the production and decay properties of the X are determined by distances at least as short as the confinement range and are sensitive to the dynamics in the core. These properties include in particular the production of the $X(3872)$ in B decays and in hadronic collisions, as well as the decays of X into final states with charmonium resonances, and into light hadrons. Given the present state of understanding the strong dynamics in the confinement region, it appears that any guidelines from the general properties of QCD, even quite approximate, may serve as a useful starting point in the studies of these properties.

The purpose of this paper is to point out a simple approximate spin selection rule for the charmed quark-antiquark pair in $X(3872)$ considered as a $J^{PC} = 1^{++}$ 'molecular' system in the previously described sense. The parameter for applicability of this rule is Λ_{QCD}/m_c. Although for the realistic charmed quark mass this parameter does not look reliably small, it still might be helpful for understanding the processes involving $X(3872)$. Namely, it will be argued here that up to corrections of order Λ_{QCD}/m_c, the Fock sum in eq.(1) should contain only the states where the total spin $S_{c\bar{c}}$ of the $c\bar{c}$ quark pair is equal to one: $S_{c\bar{c}} = 1$, irrespectively of the angular momentum and/or the overall color of this pair. Conversely, the coefficients of the states with $S_{c\bar{c}} = 0$ in the Fock sum should be suppressed by the "small" parameter[1].

As an illustration the decays $X \to \pi^0\,\chi_{cJ}$ will be discussed here, for which the spin selection rule determines the relative rate of the transitions to the charmonium states with different J.

The argumentation for the discussed here selection rule can be started with noticing that in the S-wave C-even meson pair $(D^0\bar{D}^{*0}+\bar{D}^0 D^{*0})$ the $c\bar{c}$ quark pair is necessarily in the pure spin state with $S_{c\bar{c}} = 1$. Indeed, if such meson pair is considered as a four-quark state $c\bar{c}\,u\bar{u}$

[1] The coefficients of the states without the hidden charm, i.e. only with light quarks/hadrons are trivially suppressed by m_c^{-1}, since the charmed quarks have to annihilate at distances $\sim m_c^{-1}$.

with all the relative orbital angular momenta equal to zero, the total angular momentum $J = 1$ is determined as the vector sum of the spins of the quarks and antiquarks. Each of the quark pairs, $c\bar{c}$ and $u\bar{u}$ generally can have either $S = 1$ or $S = 0$. However, one can readily see that the states where $J = 1$ arises from combining $S = 1$ from one pair with $S = 0$ of the other pair have negative C parity. The only combination resulting in a C-even $J = 1$ state is where each quark-antiquark pair has $S = 1$. Naturally, the total spin of the light quark pair is not "traceable" since it is changed by the strong interaction with amplitude of order one. Indeed, the energy gap between the states with different $S_{u\bar{u}}$ is of order Λ_{QCD} and the strong mixing amplitude is of the same order. The situation however is quite different for the total spin of the heavy quark pair. The spin-flip amplitude for a heavy charmed quark contains the factor m_c^{-1}, while the energy gap is still of order Λ_{QCD}. It should be emphasized that the latter gap arises not from the spin-dependent interactions within the $c\bar{c}$ pair (which interactions have another extra factor of m_c^{-1}), but rather from the rearrangement of the spin and/or orbital state of the light quark pair and/or changing the number of 'valence' gluons in the wave function, imposed by the conservation of the total angular momentum of the system and its P and C parities. Thus the sates ψ_i in the sum in eq.(1), resulting from the strong-interaction mixing with the 'molecular' meson state, should obey the stated spin selection rule.

Applying the spin selection rule to hadronic transitions from the $X(3872)$ to charmonium levels, one readily concludes that such transitions to the spin-singlet para-charmonium levels, e.g. $X \to \pi\,\pi\,\eta_c$ should be suppressed, while those to the spin-triplet ones should be favored. The latter include the observed decays $X(3872) \to \pi^+\,\pi^-\,J/\psi$ and $X(3872) \to \pi^+\,\pi^-\,\pi^0\,J/\psi$ and also the yet unobserved (and apparently not yet mentioned in the literature) decays to the spin-triplet P-wave states χ_{cJ}. The kinematics and the quantum numbers allow such decays with a P-wave emission of a single π^0: $X \to \pi^0\,\chi_{cJ}$. The discussed here conservation of the total spin of the heavy quark-antiquark pair implies the relation between the rates of these decays into states with different J:

$$\Gamma(X \to \pi^0\,\chi_{cJ}) \propto (2J+1)\,p_\pi^3\ , \tag{2}$$

where p_π is the momentum of the pion.

Understandably, there is a great uncertainty in estimating the absolute rate of the decays $X \to \pi^0\,\chi_{cJ}$. It is clear that as compared to the decay $X \to \pi^+\,\pi^-\,J\psi$, viewed as $X \to \rho J/\psi$, the amplitude of these decays should contain a dimensional factor μ^{-1} describing the

excitation of the heavy quark P-wave, which should naturally be of order of the characteristic size of the component of the core with the $c\bar{c}$ pair being in P wave. In terms of this factor the ratio of the decay rates can be estimated as

$$\frac{\Gamma(X \to \pi^0\, \chi_{cJ})}{\Gamma(X \to \pi^+ \pi^- J\psi)} = \frac{2J+1}{9} \, \frac{p_\pi^3}{(q_\rho)_{eff}\, \mu^2} ~, \tag{3}$$

where $(q_\rho)_{eff}$ is the effective momentum of the ρ meson in the decay $X \to \pi^+ \pi^- J\psi$ reflecting the fact that the ρ^0 materializes in the process as two pions rather than as a single particle. The value of $(q_\rho)_{eff}$ is found as

$$(q_\rho)_{eff} = \int_{4m_\pi^2}^{\Delta^2} \sqrt{\Delta^2 - q^2} \, \frac{m_\rho\, \Gamma_\rho(q^2)}{(q^2 - m_\rho^2)^2 + m_\rho^2\, \Gamma_\rho^2(q^2)} \, \frac{dq^2}{\pi} \approx 120\, MeV ~, \tag{4}$$

where $\Delta = M(X) - m_\rho$, and $\Gamma_\rho(q^2)$ is the width parameter of the ρ meson with $\Gamma_\rho(m_\rho^2) = \Gamma_\rho$. Using then eq.(3) for an estimate of the decay with e.g. the χ_{c1} in the final state (the most advantageous from the point of tagging through $\chi_{c1} \to \gamma\, J/\psi$), one finds

$$\frac{\Gamma(X \to \pi^0\, \chi_{c1})}{\Gamma(X \to \pi^+ \pi^- J\psi)} \approx 0.35 \left(\frac{0.5\, GeV}{\mu} \right)^2 ~, \tag{5}$$

which shows that the discussed decays should have a realistically observable rate for reasonable values of the parameter μ.

Considering the decays of the $X(3872)$ into light hadrons, where the $c\bar{c}$ quark pair has to annihilate, one can also use the suggested spin selection rule in combination with the "charm burning" mechanism[6], according to which the charmed quarks can annihilate from a 'molecular' state being not necessarily in a color-singlet state of the $c\bar{c}$ pair. The spin selection rule however requires the annihilation to proceed from a spin-triplet state. Then the annihilation rate is the largest from the color-octet $^3S_1$ state and is determined by the annihilation $c\bar{c} \to q\bar{q}$ through one gluon, since such state does not annihilate into two gluons in the lowest order in α_s[15, 16]. The rate of such decay is determined by the relevant size parameter of the core $(\mu')^{-1}$ and by the probability weight factor of the core:

$$\Gamma(X \to \text{light hadrons}) \sim |a_{core}|^2 \alpha_s^2(m_c) \frac{(\mu')^3}{m_c^2} ~. \tag{6}$$

Given that the core probability weight is perceived as "few percent" and that the $O(\alpha_s^2)$ annihilation rate of the charmonium states is "few MeV", the best estimate of the factors in this relation can at present be formulated only as "few percent of few MeV", i.e. in

the range from about a hundred to few hundred KeV. As uncertain as such guesstimate of the annihilation rate is, it is at least not in an apparent disagreement with the known experimental facts about the $X(3872)$.

The discussed spin selection rule might also be helpful in understanding the production of the $X(3872)$ in hadronic processes. Indeed, the production of this resonance through its peripheral 'molecular' component, i.e. by coalescence of the charmed mesons is extremely weak[17], so that the actual processes proceeds through production of the core component of the X. It can be remarked that (irrespectively of the discussed spin selection rule) this picture agrees, at least semi-quantitatively, with the rate[2] of the observed decays of $B \to K\, X(3872)$:

$$\frac{\mathcal{B}(B^+ \to K^+\, X)\, \mathcal{B}(X \to \pi^+\, \pi^-\, J/\psi)}{\mathcal{B}(B^+ \to K^+\, \psi')\, \mathcal{B}(\psi' \to \pi^+\, \pi^-\, J/\psi)} = 0.063 \pm 0.014 \; . \tag{7}$$

Indeed, it is an experimental fact that the known charmonium states are produced in B decays in association with a single Kaon with approximately the same rate (within a factor of two). One might expect then that the core states of the $X(3872)$ are produced in similar decays at approximately the same rate, so that the only suppression factor is that of the probability weight $|a_{core}|^2$, i.e. few percent. Thus there is no dramatic disagreement with the experimental number (7), if one reasonably assumes that the branching ratio $\mathcal{B}(X \to \pi^+\, \pi^-\, J/\psi)$ is not too small in comparison with $\mathcal{B}(\psi' \to \pi^+\, \pi^-\, J/\psi) \approx 0.3$.

The spin selection rule obviously somewhat restricts the possible states of the $c\bar{c}$ pair produced in a hadronic process, which fragment into the core of the $X(3872)$. It should be noted that this rule is less restrictive than recently suggested by Braaten[18], where only the production of a color-singlet ${}^3P_1$ and a color-octet ${}^3S_1$ $c\bar{c}$ states are considered (with further assumptions about the relative contributions of these mechanisms). The discussed here spin selection rule generally allows production of the $X(3872)$ originating from other states of $c\bar{c}$ as long as those states are spin triplets. It is a matter of further study whether the more restrictive assumptions of Ref.[18] are applicable in the actual production processes.

In summary. The interpretation of the resonance $X(3872)$ as a $J^{PC} = 1^{++}$ 'molecular' state of neutral D and D^* mesons is gaining support from experimental data. It is argued here that among the configurations present in the wave function describing the structure of such state in terms of quarks and gluons, should dominate those where the total spin of the $c\bar{c}$ pair is equal to one: $S_{c\bar{c}} = 1$. The parameter for suppression of the spin-singlet configurations is Λ_{QCD}/m_c. In particular this rule predicts suppression of decays of the X resonance into spin-singlet charmonium states, such as e.g. $X \to \pi\, \pi\, \eta_c$, in comparison with its hadronic

transitions to spin-triplet charmonium. The latter transitions include the observed decays with the J/ψ resonance in the final state and also the yet unobserved decays $X \to \pi^0 \, \chi_{cJ}$, for which the spin selection rule determines the ratio of the rates, and their absolute rate is likely to be within the reach of experiment. The spin selection rule favors the pattern of the decay of the X resonance into light hadrons determined by the annihilation of the $c\bar{c}$ pair into light quarks through one gluon: $c\bar{c} \to q\bar{q}$, and also the same rule can be helpful in the studies of the production of $X(3872)$ in hadronic processes.

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References

[1] K. Abe *et.al.*, (Belle Coll), Report BELLE-CONF-0439, Aug. 2004; [hep-ex/0408116].

[2] S.-K. Choi *et.al.*, (Belle Coll), Phys.Rev.Lett. **91**, 262001 (2003).

[3] D. Acosta *et.al.*, (CDF II Coll), Report FERMILAB-PUB-03-393-E, Dec 2003; [hep-ex/0312021]

[4] V.M. Abazov *et.al.*, (D0 Coll), Report FERMILAB-PUB-04-061-E, May 2004; [hep-ex/0405004]

[5] B. Aubert *et.al.*, (BABAR Coll), Report SLAC-PUB-10475, June 2004; [hep-ex/0406022]

[6] M.B. Voloshin and L.B. Okun, JETP Lett. **23**, 333 (1976).

[7] A. de Rujula, H. Georgi and S.L. Glashow, Phys.Rev.Lett. **38**, 317 (1977).

[8] N.A. Törnqvist, Phys.Rev.Lett. **67**, 556 (1991).

[9] T.E.O. Ericson and G. Karl, Phys.Lett.B **309**, 426 (1993).

[10] F.E. Close and P.R. Page, Phys.Lett.B **578**, 119 (2004).

[11] S. Pakvasa and M. Suzuki, Phys.Lett.B. **579**, 67 (2004).

[12] M.B. Voloshin, Phys.Lett.B. **579**, 316 (2004).

[13] N.A. Törnqvist, Phys.Lett.B **590**, 209 (2004); [hep-ph/0402237] and [hep-ph/0308277].

[14] E.S. Swanson, Phys.Lett.B **588**, 189 (2004).

[15] R. Barbieri, R. Gatto, and E. Remiddi, Phys.Lett.B **61**, 465 (1976).

[16] L.B. Okun and M.B. Voloshin, Report ITEP-152 (1976), unpublished, included into V.A. Novikov *et.al.*, Phys.Rep.C **41**, 1 (1978).

[17] E. Braaten and M. Kusunoki, Phys.Rev.D **69**, 114012 (2004).

[18] E. Braaten, Ohio State Univ. Report hep-ph/0408230, Aug. 2004.